

Natural μ -term with Peccei-Quinn Symmetry

E. J. Chun*

International Centre for Theoretical Physics
P.O.Box 586, 34100 Trieste, Italy

Abstract

The generalized Higgs mass term NH_1H_2 of the supersymmetric standard model is used to implement the Peccei-Quinn Symmetry to solve the strong-CP problem. Then supersymmetry breaking can generate the Higgs mass parameter μ of order $m_{3/2}$ through soft-breaking parameters. This kind of extension contains extra light fields of the axion supermultiplet whose dominant coupling may come from the supersymmetric axion-Higgs-Higgs coupling. We present a working example and discuss the cosmological implications of the model.

*Email address: chun@ictp.trieste.it

The supersymmetric Higgs mass term $\mu H_1 H_2$ in the minimal supersymmetric standard model (MSSM) brings a problem of naturalness, so called the μ -problem [1]. The μ -term together with soft supersymmetry breaking terms drive electroweak symmetry breaking. Therefore both the parameter μ and soft-breaking parameters are required to be at or slightly above the electroweak scale. The scale of soft-breaking parameters (characterized by the gravitino mass $m_{3/2}$) can be understood from the hidden-sector supersymmetry breaking mechanism [2]. The question is why the μ -parameter is so small compared to the other scale in the theory, e.g., the Planck scale M_{Pl} . The μ -term may well be generated dynamically due to supersymmetry breaking.

One conventional way to explain the origin of the μ -term is to include a singlet N under the standard model (SM) gauge group and introduce a term $NH_1 H_2$ [3]. Then one can arrange for the vacuum expectation value (VEV) of N to get the proper value. In this case one generically expects the presence of extra light singlets. Very recently it has been demonstrated that introducing an $U(1)$ gauge group can achieve the generation of the μ -term without producing light singlets [4]. This approach can make N as heavy as M_{Pl} but calls for some light colored fields in order to cancel the anomaly.

In this letter we stress that the generation of the μ -term can be made much more appealing if one uses the Peccei-Quinn (PQ) Symmetry [5] instead of any other global or local symmetry. Obviously one can combine the μ -problem and the strong-CP problem [6] in this way. If one wants to understand the strong-CP problem in terms of the PQ mechanism, it is quite natural to assign the presence of the term $hNH_1 H_2$ to the PQ symmetry. Then one has to arrange certain Higgs superpotential providing PQ symmetry breaking at the invisible scale $M_{PQ} \sim 10^{10} - 10^{12}$ GeV. In this case $\langle N \rangle \sim M_{PQ}$ implies extreme fine-tuning of h . But it is not necessary for N to have such a large VEV as is commonly believed. If the VEV of N vanishes in the supersymmetric limit, soft supersymmetry breaking may induce nonvanishing VEV which is expected to be of order $m_{3/2}$. The PQ symmetry is broken by some other fields which couple to N .

In the first attempt to relate the dynamical generation of the μ -term to the strong-CP problem, a non-renormalizable term like $S^2 H_1 H_2 / M_{Pl}$ was used [1]. For this to provide the proper value of μ , the PQ scale ($\langle S \rangle \sim M_{PQ}$) is necessarily of the same order as the hidden sector supersymmetry breaking scale which is indeed inside the above-mentioned

window for the PQ scale. However our prescription with the renormalizable term is in fact irrelevant to the size of the PQ scale, which makes it viable even if any dissipation mechanism of the axion energy density works to remove the upper bound of the PQ scale [7].

Other approaches to the μ -problem in supergravity or superstring theories have been investigated in refs. [8] with the PQ symmetry and in refs. [9] and [10] without it.

In the below we will illustrate that the VEV of N is generated due to soft-breaking parameters and the PQ symmetry breaking is driven by the VEV's of some other fields (called S). Only one light degree of freedom (that is, the axion supermultiplet Φ) is added to the MSSM. The couplings of the axion supermultiplet to the MSSM particles are well-fixed due to the nature of the PQ symmetry. The generalized μ -term NH_1H_2 induces the supersymmetric axion-Higgs-Higgs vertex ΦH_1H_2 which may provide the dominant coupling of the axion and its superpartners. The mass splitting inside the axion supermultiplet and its cosmological role through the μ -term coupling will be also discussed.

Our working example introduces five ‘‘PQ fields’’ N, S'', S, S' and Y which are singlets under the SM gauge group and their PQ charges are assigned to be $X = (-2, 2, 1, -1, 0)$. Furthermore we impose the R-symmetry under which the PQ fields carry the charges $R = (2, 0, 0, 0, 2)$. The most general ‘‘PQ superpotential’’ invariant under both the PQ symmetry and the R-symmetry is

$$W_{PQ} = MNS'' + fNS^2 + f'(SS' - M'^2)Y \quad (1)$$

together with the term hNH_1H_2 . Note that the renormalizable term $S''S'^2$ and the non-renormalizable term $H_1H_2S'^2/M_{Pl}$ are forbidden by the R-symmetry, which is crucial for our discussion. In the supersymmetric limit, the VEV's are given by

$$\begin{aligned} \langle SS' \rangle &= M'^2 \\ \langle S'' \rangle &= -f\langle S \rangle^2/M \\ \langle Y \rangle &= 0 \\ \langle N \rangle &= 0. \end{aligned} \quad (2)$$

The Goldstone (axion) mode is a linear combination of S, S' and S'' . The PQ scale is given by $M_{PQ} = (\langle S \rangle^2 + \langle S' \rangle^2 + 4\langle S'' \rangle^2)^{1/2}$. One can see that all the other modes get masses of order M_{PQ} .

Inclusion of soft supersymmetry breaking effect gives the desired feature. The scalar potential should be completed with the soft-terms,

$$V_{soft} = BMNS'' + fANS^2 + f'A'SS'Y - f'B'M'^2Y + \text{h.c.} \\ + m_N^2|N|^2 + m_S^2|S|^2 + m_{S'}^2|S'|^2 + m_{S''}^2|S''|^2 \quad (3)$$

where A, B, A', B' and m_i are soft-breaking parameters of order $m_{3/2}$. Minimization of the full scalar potential yields the following changes from eq. (2):

$$\langle N \rangle = \frac{2f(A' - B') - f(1 + r^2)(A - B)}{\xi(r + r^{-1}) + 4f^2r^2} \quad (4) \\ \langle Y \rangle = \frac{2f^2r^{-1}(A - B) - f'^{-1}(\xi + 4f^2r^{-1})(A' - B')}{\xi(r + r^{-1}) + 4f^2r^2}$$

where $r \equiv \langle S/S' \rangle$ and $\xi \equiv M^2/M'^2$. The value of r depends on the parameters, which we do not show explicitly. The VEV's of S, S' and S'' get negligible changes of order $m_{3/2}^2/M_{Pl}$. We note that in the limit $f \rightarrow 0$, we have $\langle N \rangle \rightarrow 0$ and $\langle Y \rangle \rightarrow -(A' - B')/f'$ as discussed in the connection with supersymmetric majoron models [11].

We have seen that the light degrees of freedom contain the axion supermultiplet in addition to the usual MSSM particles. The couplings of the axion supermultiplet to the MSSM particles are just supersymmetric counter parts of the conventional Dine-Fischler-Srednicki-Zhitnitskii axion model [12]. In addition, the generalized μ -term hNH_1H_2 implies the supersymmetric axion-Higgs-Higgs coupling

$$W_{ahh} = X_N \frac{\mu}{M_{PQ}} \Phi H_1 H_2 \quad (5)$$

where $X_N = -2$ and $\mu = h\langle N \rangle$. It can play an important role in the supersymmetric axion models. The mass splitting among the axion and its supersymmetric particles (called the saxion and the axino) can be also calculated in our model. The axino mass is given by $m_{\tilde{a}} = f'\langle Y \rangle$ and the saxion mass squared is certain linear combination of $m_S^2, m_{S'}^2$ and $m_{S''}^2$. Both of them get masses of order $m_{3/2}$ without accepting any fine-tuning of the parameters.

We now turn to the cosmology of the model. Since the parameter μ is of order 100 GeV, the above axion-Higgs-Higgs vertex may give the most strongest interaction of the axion supermultiplet to the MSSM fields. The axino and the saxion with masses of order $m_{3/2}$ have to be unstable [13]. Then the axino (saxion) may decay into a Higgsino and a Higgs (two Higgses). A cosmological bound on the axino and the saxion mass comes from the standard nucleosynthesis. Since the axino and the saxion decouple at very high temperature, their relic densities overdominate the energy density of the universe when they decay. Then in order not to destroy the prediction of the nucleosynthesis their lifetime should be shorter than about 1 sec. The axion-Higgs-Higgs vertex induces fast enough decay to satisfy this constraint. For the axino we have another constraint from the fact that its decay products should contain at least one lightest supersymmetric particle (LSP) which forms cold dark matter of the universe. In order to avoid the overclosure due to the decay-produced LSP's, the axino decay should occur before the LSP decouples. Taking the typical decoupling temperature of the LSP at about $T = m_{LSP}/20$, we get the bound on the axino mass

$$m_{\tilde{a}} \gtrsim 190 \text{ GeV} \left(\frac{m_{LSP}}{20 \text{ GeV}} \right)^2 \left(\frac{200 \text{ GeV}}{\mu} \right)^2 \left(\frac{M_{PQ}}{10^{12} \text{ GeV}} \right)^2. \quad (6)$$

If μ is smaller than the top quark mass, the axino may decay into a top and a light stop in which case one replace μ by the top quark mass.

The axino itself can be the LSP when it is lighter than a few keV providing the underclosure energy density of the universe. One may be able to find a PQ superpotential which admits such a light axino [14].

If one would like to have the μ -parameter smaller than 100 GeV or the larger m_{LSP} , the above bound produces a bit large amount of splitting between the axino mass and the μ -parameter. In the presented model, the required splitting can be obtained by tuning the parameters f and f' . However, better way to overcome this constraint is to invoke inflation. The decoupling temperature of the axino (saxion) may be higher than the reheating temperature $T_R \sim 10^{10}$ GeV which is a maximally allowed value to cure the gravitino problem in supergravity models [15]. Then the primordial axino relics are diluted away. We recall that the axino decoupling temperature is determined by the annihilation

of an axino and a gluino into two quarks [13], which gives the decoupling temperature;

$$T_D \sim 10^{11} \text{ GeV} \left(\frac{M_{PQ}}{10^{12} \text{ GeV}} \right)^2 \left(\frac{0.1}{\alpha_c} \right)^3. \quad (7)$$

After inflation, the universe can produce the axino population through inequilibrium process of the inverse annihilation. The ratio of the regenerated number density to the entropy density is given by [15]

$$Y \sim 7 \times 10^{-5} \left(\frac{10^{12} \text{ GeV}}{M_{PQ}} \right)^2 \left(\frac{T_R}{10^{10} \text{ GeV}} \right). \quad (8)$$

The regenerated number of axinos should be sufficiently reduced in order to avoid overclosure due to the decay-products. It gives the upper bound on the reheating temperature;

$$T_R \lesssim 1.4 \times 10^7 \text{ GeV} \left(\frac{M_{PQ}}{10^{12} \text{ GeV}} \right)^2 \left(\frac{20 \text{ GeV}}{m_{LSP}} \right). \quad (9)$$

Therefore, either the lower bound on the axino mass (6) or the upper bound on the reheating temperature (9) has to be satisfied in the inflationary universe.

In conclusion, we have illustrated a mechanism of generating the μ -term in the axionic extension of the MSSM. Soft supersymmetry breaking parameters induce the natural value of order $m_{3/2}$ for the parameter μ . This scheme is precisely the supersymmetric version of the Dine-Fischler-Srednicki-Zhitnitskii axion model which adds the light fields from the axion supermultiplet to the MSSM. The supersymmetrized axion-Higgs-Higgs interaction induced from the extended Higgs mass term NH_1H_2 can provide the main decay mode of the superpartners of the axion. From this we have drawn the bound on the axino mass or on the reheating temperature in the inflationary universe.

Acknowledgement: The author would like to thank Professor Abdus Salam, the International Atomic Energy Agency and UNESCO for hospitality at the International Centre for Theoretical Physics, Trieste.

References

- [1] J.E. Kim and H.P. Nilles, Phys. Lett. B138 (1984) 150.
- [2] H.P. Nilles, Phys. Rep. 110 (1984) 1.
- [3] H.P. Nilles, M. Srednicki and D. Wyler, Phys. Lett. B120 (1983) 346; J.M. Frere, D.R.T. Jones and S. Raby, Nucl. Phys. B222 (1983)11; J.P. Derendinger and C. Savoy, Nucl. Phys. B237 (1984) 307; L.E. Ibanez and J. Mas, Nucl. Phys. B286 (1987) 107; J. Ellis, J.F. Gunion, H.E. Haber, L. Roszkowski and F. Zwirner, Phys. Rev. D39 (1989) 844; M. Drees, Int. J. Mod. Phys. A4 (1993) 288.
- [4] R. Hempfling, Phys. Lett. B329 (1994) 222.
- [5] R.D. Peccei and H.R. Quinn, Phys. Rev. Lett. 38 (1977) 1440; R.D. Peccei and H.R. Quinn, Phys. Rev. D16 (1977) 1791.
- [6] For reviews and more references see, J.E. Kim, Phys. Rep. 150 (1987) 1; H.-Y. Cheng, Phys. Rep. 158 (1988) 1; R.D. Peccei, in *CP Violation*, ed. C. Jarlskog (WSPC, Singapore, 1989) 503.
- [7] K.S. Babu, S.M. Barr and D. Seckel, preprint BA-94-21, hep-ph/9406308.
- [8] E.J. Chun, J.E. Kim and H.P. Nilles, Nucl. Phys. B370 (1992) 105; J.E. Kim and H.P. Nilles, preprint SNUTP 94-55, hep-ph/9406296.
- [9] G.F. Giudice and A. Masiero, Phys. Lett. B206 (1988) 480; J.A. Casas and C. Munoz, Phys. Lett. B306 (1993) 288; A. Brignole, L.E. Ibanez and C. Munoz, Nucl. Phys. B422 (1994) 125; I. Antoniadis, E. Gava, K.S. Narain and T.R. Taylor, preprint IC/94/72 and hep-th/9405254.
- [10] S. Ferrara, C. Kounnas, M. Porrati and F. Zwirner, Nucl. Phys. B318 (1989) 75; S. Ferrara, C. Kounnas and F. Zwirner, Nucl. Phys. B429 (1994) 589; A. Brignole and F. Zwirner, preprint CERN-TH.7439/94, hep-th/9409099 (to appear in Phys. Lett. B).

- [11] R.N. Mohapatra and X. Zhang, Phys. Rev. D49 (1994) R1163; E.J. Chun, H.B. Kim and A. Lukas, Phys. Lett. B328 (1994) 346.
- [12] A.R. Zhitnitskii, Sov. J. Nucl. Phys. 31 (1980) 103; M. Dine, W. Fischler and M. Srednicki, Phys. Lett. B104 (1981) 199.
- [13] K. Rajagopal, M.S. Turner and F. Wilczek, Nucl. Phys. B358 (1994) 447.
- [14] T. Goto and M. Yamaguchi, Phys. Lett. B276 (1992) 103; E.J. Chun, J.E. Kim and H.P. Nilles, Phys. Lett. B287 (1992) 123.
- [15] J. Ellis, J. E. Kim and D. V. Nanopoulos, Phys. Lett. B145 (1984) 181; R. Juskiwicz, J. Silk and A. Stebbins, Phys. Lett. B158 (1985) 463; J. Ellis, D. V. Nanopoulos and S. Sarkar, Nucl. Phys. B219 (1985) 175; M. Kawasaki and K. Sato, Phys. Lett. B189 (1987) 23; T. Moroi, and H. Murayama and M. Yamaguchi, Phys. Lett. B303 (1993) 289.